

E_6 inspired composite Higgs model and 750 GeV diphoton excess

Roman Nevzorov^{1,2,*} and Anthony Thomas¹

¹ARC Centre of Excellence for Particle Physics at the Terascale and CSSM, Department of Physics, The University of Adelaide, Adelaide SA 5005, Australia

²Institute for Theoretical and Experimental Physics, Moscow 117218, Russia

Abstract. In the E_6 inspired composite Higgs model (E_6 CHM) the strongly interacting sector possesses an $SU(6) \times U(1)_B \times U(1)_L$ global symmetry. Near scale $f \gtrsim 10$ TeV the $SU(6)$ symmetry is broken down to its $SU(5)$ subgroup, that involves the standard model (SM) gauge group. This breakdown of $SU(6)$ leads to a set of pseudo–Nambu–Goldstone bosons (pNGBs) including a SM–like Higgs and a SM singlet pseudoscalar A . Because of the interactions between A and exotic fermions, which ensures the approximate unification of the SM gauge couplings and anomaly cancellation in this model, the couplings of the pseudoscalar A to gauge bosons get induced. As a result, the SM singlet pNGB state A with mass around 750 GeV may give rise to sufficiently large cross section of $pp \rightarrow \gamma\gamma$ that can be identified with the recently observed diphoton excess.

1 Introduction

The discovery of the Higgs boson with mass $m_h \simeq 125$ GeV allows one to estimate the values of parameters of the Higgs potential. In the standard model (SM) the Higgs scalar potential is given by

$$V(H) = m_H^2 H^\dagger H + \lambda (H^\dagger H)^2. \quad (1)$$

The 125 GeV Higgs mass corresponds to $m_H^2 \approx -(90 \text{ GeV})^2$ and $\lambda \approx 0.13$. At the moment current data does not permit to distinguish whether Higgs boson is an elementary particle or a composite state. Although the discovered Higgs boson can be composed of more fundamental degrees of freedom, the rather small values of $|m_H^2|$ and the Higgs quartic coupling λ indicate that the Higgs field may emerge as a pseudo–Nambu–Goldstone boson (pNGB) from the spontaneous breaking of an approximate global symmetry of some strongly interacting sector.

The minimal composite Higgs model (MCHM) [1] includes weakly–coupled elementary and strongly coupled composite sectors (for a recent review, see [2]). The weakly–coupled elementary sector involves all SM fermions and gauge bosons. The strongly coupled sector gives rise to a set of bound states that contains Higgs doublet and massive fields with the quantum numbers of all SM particles. These fields are associated with the composite partners of the quarks, leptons and gauge bosons. The elementary states couple to the composite operators of the strongly interacting sector leading to

*e-mail: roman.nevzorov@adelaide.edu.au

mixing between these states and their composite partners. In this framework, which is called partial compositeness, the couplings of the SM states to the composite Higgs are set by the fractions of the compositeness of these states. The observed mass hierarchy in the quark and lepton sectors can be accommodated through partial compositeness if the fractions of compositeness of the first and second generation fermions are quite small. In this case the flavour-changing processes and the modifications of the W and Z couplings associated with the light SM fermions are somewhat suppressed. At the same time, the top quark is so heavy that the right-handed top quark t^c should have sizeable fraction of compositeness.

The strongly interacting sector of the MCHM possesses global $SO(5) \times U(1)_X$ symmetry that contains the $SU(2)_W \times U(1)_Y$ subgroup. Near the scale f the $SO(5)$ symmetry is broken down to $SO(4)$ so that the SM gauge group remains intact, resulting in four pNGB states which form the Higgs doublet. The custodial global symmetry $SU(2)_{cust} \subset SO(4)$ allows one to protect the Peskin–Takeuchi \hat{T} parameter against new physics contributions. Experimental limits on the parameter S imply that $m_\rho = g_\rho f \gtrsim 2.5$ TeV, where m_ρ is a scale associated with the masses of the set of spin-1 resonances and g_ρ is a coupling of these ρ -like vector resonances. This set of resonances, in particular, contains composite partners of the SM gauge bosons. Even more stringent bounds on f come from the non-observation of flavor changing neutral currents (FCNCs). In the composite Higgs models, adequate suppression of the non-diagonal flavour transitions can be obtained only if f is larger than 10 TeV. This bound on the scale f can be considerably alleviated in the models with additional flavour symmetries FS. In the models with $FS = U(2)^3 = U(2)_q \times U(2)_u \times U(2)_d$ symmetry, the bounds that originate from the Kaon and B systems can be satisfied even for $m_\rho \sim 3$ TeV. In these models the appropriate suppression of the baryon number violating operators and the Majorana masses of the left-handed neutrino can be achieved if global $U(1)_B$ and $U(1)_L$ symmetries, which ensure the conservation of the baryon and lepton numbers to a very good approximation, are imposed. Thus the composite Higgs models under consideration are based on

$$SU(3)_C \times SO(5) \times U(1)_X \times U(1)_B \times U(1)_L \times FS. \quad (2)$$

The couplings of the elementary states to the strongly interacting sector explicitly break the $SO(5)$ global symmetry. As a consequence, the pNGB Higgs potential arises from loops involving elementary states. This leads to the suppression of the effective quartic Higgs coupling λ .

2 E_6 inspired composite Higgs model

In the E_6 inspired composite Higgs model (E_6 CHM) the Lagrangian of the strongly coupled sector is invariant under the transformations of an $SU(6) \times U(1)_B \times U(1)_L$ global symmetry. The E_6 CHM can be embedded into $N = 1$ supersymmetric (SUSY) orbifold Grand Unified Theories (GUTs) in six dimensions which are based on the $E_6 \times G_0$ gauge group [3]. (Different aspects of the E_6 inspired models with low-scale supersymmetry breaking were recently considered in [4]–[19].) Near some high energy scale, M_X , the $E_6 \times G_0$ gauge group is broken down to the $SU(3)_C \times SU(2)_W \times U(1)_Y \times G$ subgroup where $SU(3)_C \times SU(2)_W \times U(1)_Y$ is the SM gauge group. Gauge groups G_0 and G are associated with the strongly interacting sector. Fields belonging to this sector can be charged under both the E_6 and G_0 (G) gauge symmetries. The weakly-coupled sector includes elementary states that participate in the E_6 interactions only. Due to the conservation of the $U(1)_B$ and $U(1)_L$ charges all elementary states with different baryon and/or lepton numbers are components of different bulk 27-plets, whereas all other components of these 27-plets have to acquire masses of the order of M_X . All fields from the strongly interacting sector reside on the brane where E_6 symmetry is broken down to the $SU(6)$ that contains an $SU(3)_C \times SU(2)_W \times U(1)_Y$ subgroup. As a result at high energies the

Lagrangian of the composite sector respects SU(6) global symmetry. The SM gauge interactions violate this symmetry. Nevertheless, SU(6) can remain an approximate global symmetry of the strongly coupled sector at low energies if the gauge couplings of this sector are considerably larger than the SM ones.

As in most composite Higgs models, the global SU(6) symmetry in the E₆CHM is expected to be broken below scale f . Here we assume that it gets broken to SU(5) subgroup, so that the SM gauge group is preserved. Since E₆CHM does not possess any extra custodial or flavour symmetry, the scale f must be much larger than the weak scale, i.e. $v \ll f$. In particular, the adequate suppression of the FCNCs requires $f \gtrsim 10$ TeV. The SU(6)/SU(5) coset space includes eleven pNGB states that correspond to the broken generators $T^{\hat{a}}$ of SU(6). These pNGB states can be parameterised by

$$\begin{aligned} \Omega^T &= \Omega_0^T \Sigma^T = e^{i \frac{\phi_0}{\sqrt{15}f}} \begin{pmatrix} C\phi_1 & C\phi_2 & C\phi_3 & C\phi_4 & C\phi_5 & \cos \frac{\tilde{\phi}}{\sqrt{2}f} + \sqrt{\frac{3}{10}} C\phi_0 \end{pmatrix}, \\ C &= \frac{i}{\tilde{\phi}} \sin \frac{\tilde{\phi}}{\sqrt{2}f}, \quad \tilde{\phi} = \sqrt{\frac{3}{10} \phi_0^2 + |\phi_1|^2 + |\phi_2|^2 + |\phi_3|^2 + |\phi_4|^2 + |\phi_5|^2}, \end{aligned} \quad (3)$$

where the SU(6) generators are normalised so that $\text{Tr} T^a T^b = \frac{1}{2} \delta_{ab}$ and

$$\Omega_0^T = (0 \quad 0 \quad 0 \quad 0 \quad 0 \quad 1), \quad \Sigma = e^{i\Pi/f}, \quad \Pi = \Pi^{\hat{a}} T^{\hat{a}}.$$

In the leading approximation the Lagrangian, that describes the interactions of the pNGB states, can be written as

$$\mathcal{L}_{\text{pNGB}} = \frac{f^2}{2} \left| \mathcal{D}_\mu \Omega \right|^2. \quad (4)$$

The field ϕ_0 is real and does not participate in the SU(3)_C × SU(2)_W × U(1)_Y gauge interactions. Five components of vector Ω , i.e. $\tilde{H} \sim (\phi_1 \ \phi_2 \ \phi_3 \ \phi_4 \ \phi_5)$, form a fundamental representation of the unbroken SU(5) subgroup of SU(6). The components $H \sim (\phi_1 \ \phi_2)$ transform as an SU(2)_W doublet. Therefore H corresponds to the SM-like Higgs doublet. Three other components of \tilde{H} , i.e. $T \sim (\phi_3 \ \phi_4 \ \phi_5)$, transform as an SU(3)_C triplet. In the E₆CHM neither H nor T carry baryon and/or lepton number.

The pNGB effective potential $V_{\text{eff}}(\tilde{H}, T, \phi_0)$ is induced by the interactions of the elementary states with their composite partners, which break SU(6) global symmetry. The analysis of the structure of this potential including the derivation of quadratic terms $m_H^2 |H|^2$ and $m_T^2 |T|^2$ in the composite Higgs models, which are similar to the E₆CHM, shows that there is a considerable part of the parameter space where m_H^2 is negative and m_T^2 is positive [20]–[21]. In this parameter region the SU(2)_W × U(1)_Y gauge symmetry gets broken to U(1)_{em}, associated with electromagnetism, while SU(3)_C colour is preserved. Because in the E₆CHM the scale $f \gtrsim 10$ TeV, a significant tuning, $\sim 0.01\%$, is required to get the appropriate value of the parameter m_H^2 that results in a 125 GeV Higgs state.

Since in the E₆CHM all states in the strongly interacting sector fill complete SU(5) representations the corresponding fields contribute equally to the beta functions of the SU(3)_C, SU(2)_W and U(1)_Y interactions in the one-loop approximation. As a consequence the convergence of the SM gauge couplings is determined by the matter content of the weakly-coupled sector. In this case, approximate gauge coupling unification can be achieved if the right-handed top quark, t^c , is entirely composite and the weakly-coupled elementary sector involves the following set of multiplets (see also [22]):

$$(q_i, d_i^c, \ell_i, e_i^c) + u_\alpha^c + \bar{q} + \bar{d}^c + \bar{\ell} + \bar{e}^c + \eta, \quad (5)$$

where $\alpha = 1, 2$ runs over the first two generations and $i = 1, 2, 3$ runs over all three. In Eq. (5) u_α^c, d_i^c and e_i^c represent the right-handed up- and down-type quarks and charged leptons, q_i and ℓ_i

correspond to the left-handed quark and lepton doublets, whereas \bar{q} , \bar{d}^c , $\bar{\ell}$ and \bar{e}^c are exotic states which have opposite $SU(3)_C \times SU(2)_W \times U(1)_Y$ quantum numbers to the left-handed quark doublets, right-handed down-type quarks, left-handed lepton doublets and right-handed charged leptons, respectively. An additional SM singlet exotic state, η , with spin 1/2 is included to ensure the phenomenological viability of the model under consideration. The set of elementary fermion states (5) is chosen so that the weakly-coupled sector contains all SM fermions except right-handed top quark and anomaly cancellation takes place. Using the one-loop renormalisation group equations (RGEs) one can find the exact gauge coupling unification is attained for $\alpha(M_Z) = 1/127.9$, $\sin^2 \theta_W = 0.231$ and $\alpha_3(M_Z) \simeq 0.109$. The scale where the unification of the SM gauge couplings takes place is somewhat close to $M_X \sim 10^{15} - 10^{16}$ GeV. This estimation demonstrates that for the phenomenologically acceptable values $\alpha_3(M_Z) \simeq 0.118$ the SM gauge couplings can be reasonably close to each other at very high energies around $M_X \simeq 10^{16}$ GeV.

The scenario under consideration implies that the dynamics of the strongly coupled sector below the scale f leads to the composite $\mathbf{10} + \mathbf{\bar{5}} + \mathbf{1}$ multiplets of $SU(5)$. Because of the conservation of the $U(1)_B$ and $U(1)_L$ charges all components of the 10-plet, i.e. Q , E^c and t^c , carry the same baryon and lepton numbers as the right-handed top quark t^c . The components of $\mathbf{\bar{5}}$ (D^c and L) and $\mathbf{1}$ ($\bar{\eta}$) can have baryon charges $-1/3$ and $+1/3$ [3]. It is expected that the composite multiplets Q , E^c , D^c , L and $\bar{\eta}$ get combined with the elementary exotic states \bar{q} , \bar{e}^c , \bar{d}^c , $\bar{\ell}$ and η , respectively, giving rise to a set of vector-like fermion states. The only exceptions are the components of the 10-plet associated with the composite right-handed top quark which survive down to the electroweak scale.

In the E_6 CHM the lightest exotic fermion state has to be stable. Indeed, the baryon number conservation implies that the Lagrangian of the E_6 CHM is also invariant under the transformations of the discrete Z_3 symmetry which can be defined as

$$\Psi \longrightarrow e^{2\pi i B_3/3} \Psi, \quad B_3 = (3B - n_C) \bmod 3. \quad (6)$$

Here B is the baryon number of the given multiplet Ψ and n_C is the number of colour indices ($n_C = 1$ for $\mathbf{3}$ and $n_C = -1$ for $\mathbf{\bar{3}}$). This symmetry is called baryon triality [20]. All states in the SM have $B_3 = 0$. At the same time exotic fermion states carry either $B_3 = 1$ or $B_3 = 2$. As a result the lightest exotic state with non-zero B_3 can not decay into SM particles and should be stable. Since models with stable charged particles are ruled out by various experiments [23]-[24], the lightest exotic fermion in the E_6 CHM must be neutral. It is also worth noting that the coupling of this neutral Dirac fermion to the Z -boson have to be extremely suppressed. Otherwise this stable exotic state would scatter on nuclei resulting in unacceptably large spin-independent cross sections. Thus, only a Dirac fermion, which is mostly a superposition of η and $\bar{\eta}$, can be the lightest exotic state in the E_6 CHM.

3 750 GeV diphoton resonance

The SM singlet pNGB state ϕ_0 can be identified with the 750 GeV diphoton resonance recently reported by ATLAS and CMS. It is important that no 750 GeV resonance has been observed in other channels like $pp \rightarrow t\bar{t}, WW, ZZ, b\bar{b}, \tau\bar{\tau}$ and jj . This may be an indication that the detected signal is just a statistical fluctuation. At the same time, if these observations are confirmed this should set stringent constraints on the new physics models that may lead to such a signature. For example, in the E_6 CHM the field ϕ_0 can mix with the Higgs boson which would result in large partial widths of the 750 GeV resonance associated with the decays of this state into pairs of Z -bosons, W -bosons and $t\bar{t}$. The corresponding mixing can be suppressed if invariance under the CP transformation is imposed. Indeed, in this case ϕ_0 manifests itself in the Yukawa interactions with fermions as a pseudoscalar

field. In particular, the couplings of the SM singlet pNGB state $\phi_0 = A$ to the top quarks is induced by

$$\mathcal{L}_{AT} = \frac{y_t}{\Lambda_t} A (i\bar{t}_L H^0 t_R + h.c.). \quad (7)$$

Because of the almost exact CP-conservation the mixing between the Higgs boson and pseudoscalar A is forbidden.

The Lagrangian that describes the interactions between A and exotic fermions can be written in the following form [25]

$$\mathcal{L}_{AF} = A \left(i\kappa_D \bar{d}^c D^c + i\kappa_Q \bar{q} Q + i\lambda_L \bar{\ell} L + i\lambda_E \bar{e}^c E^c + i\lambda_\eta \bar{\eta} \eta + h.c. \right). \quad (8)$$

In the most general case the couplings κ_i and λ_i in Eq. (8) and the exotic fermion masses μ_i , induced below scale f , i.e.

$$\mathcal{L}_{mass} = \mu_D \bar{d}^c D^c + \mu_Q \bar{q} Q + \mu_L \bar{\ell} L + \mu_E \bar{e}^c E^c + \mu_\eta \bar{\eta} \eta + h.c., \quad (9)$$

are entirely independent parameters, which are not constrained by the SU(6) and SU(5) symmetries. In order to get the cross section $\sigma(pp \rightarrow \gamma\gamma)$, which corresponds to the production and sequential diphoton decays of the pseudoscalar A , of about 5 – 10 fb we assume that $\mu_D, \mu_Q, \mu_L, \mu_E$ and μ_η are larger than 375 GeV. As a result the on-shell decays of A into the exotic fermions, that result in the strong suppression of the branching ratios of the decays of this pNGB state into photons, are not kinematically allowed. Integrating out the exotic fermion states one obtains the effective Lagrangian that describes the interactions of the pseudoscalar A with the SM gauge bosons [25]

$$\mathcal{L}_{eff}^A = c_1 A B_{\mu\nu} \tilde{B}^{\mu\nu} + c_2 A W_{\mu\nu}^a \tilde{W}^{a\mu\nu} + c_3 A G_{\mu\nu}^\sigma \tilde{G}^{\sigma\mu\nu}, \quad (10)$$

where $B_{\mu\nu}, W_{\mu\nu}^a, G_{\mu\nu}^\sigma$ are field strengths for the $U(1)_Y, SU(2)_W$ and $SU(3)_C$ gauge interactions, $\tilde{G}^{\sigma\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} G_{\lambda\rho}^\sigma, \tilde{W}^{a\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} W_{\lambda\rho}^a, \tilde{B}^{\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} B_{\lambda\rho}$, whereas

$$\begin{aligned} c_1 &= \frac{\alpha_Y}{16\pi} \left[\frac{2\kappa_D}{3\mu_D} B(x_D) + \frac{\kappa_Q}{3\mu_Q} B(x_Q) + \frac{\lambda_L}{\mu_L} B(x_L) + 2\frac{\lambda_E}{\mu_E} B(x_E) \right], \\ c_2 &= \frac{\alpha_2}{16\pi} \left[3\frac{\kappa_Q}{\mu_Q} B(x_Q) + \frac{\lambda_L}{\mu_L} B(x_L) \right], \\ c_3 &= \frac{\alpha_3}{16\pi} \left[\frac{\kappa_D}{\mu_D} B(x_D) + 2\frac{\kappa_Q}{\mu_Q} B(x_Q) \right], \\ B(x) &= 2x \arcsin^2[1/\sqrt{x}], \quad \text{for } x \geq 1. \end{aligned} \quad (11)$$

In Eq. (11) $x_D = 4\mu_D^2/m_A^2, x_Q = 4\mu_Q^2/m_A^2, x_L = 4\mu_L^2/m_A^2, x_E = 4\mu_E^2/m_A^2, m_A \simeq 750$ GeV is the mass of the SM singlet pNGB state $A, \alpha_Y = 3\alpha_1/5$ while α_1, α_2 and α_3 are (GUT normalised) gauge couplings of $U(1)_Y, SU(2)_W$ and $SU(3)_C$ interactions. Using Eqs. (10)–(11) one can obtain an analytical expression for the coupling of the pseudoscalar A to the electromagnetic field $F_{\mu\nu}$

$$\mathcal{L}_{eff}^{A\gamma\gamma} = c_\gamma A F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad c_\gamma = c_1 \cos^2 \theta_W + c_2 \sin^2 \theta_W, \quad (12)$$

where θ_W is the weak mixing (Weinberg) angle and $\tilde{F}^{\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} F_{\lambda\rho}$.

Since at the LHC the pseudoscalar A is predominantly produced through gluon fusion the cross section $\sigma_{\gamma\gamma} = \sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ can be presented in the following form [26]

$$\sigma_{\gamma\gamma} \simeq \frac{C_{gg}}{m_A s} \frac{\Gamma(A \rightarrow gg) \Gamma(A \rightarrow \gamma\gamma)}{\Gamma_A} \simeq 7.3 \text{ fb} \times \left(\frac{\Gamma(A \rightarrow gg) \Gamma(A \rightarrow \gamma\gamma)}{\Gamma_A m_A} \times 10^6 \right), \quad (13)$$

where $C_{gg} \simeq 3163$, $\sqrt{s} \simeq 13$ TeV, Γ_A is a total width of the pseudoscalar A while partial decay widths $\Gamma(A \rightarrow \gamma\gamma)$ and $\Gamma(A \rightarrow gg)$ are given by

$$\Gamma(A \rightarrow gg) = \frac{2m_A^3}{\pi} |c_3|^2, \quad \Gamma(A \rightarrow \gamma\gamma) = \frac{m_A^3}{4\pi} |c_\gamma|^2. \quad (14)$$

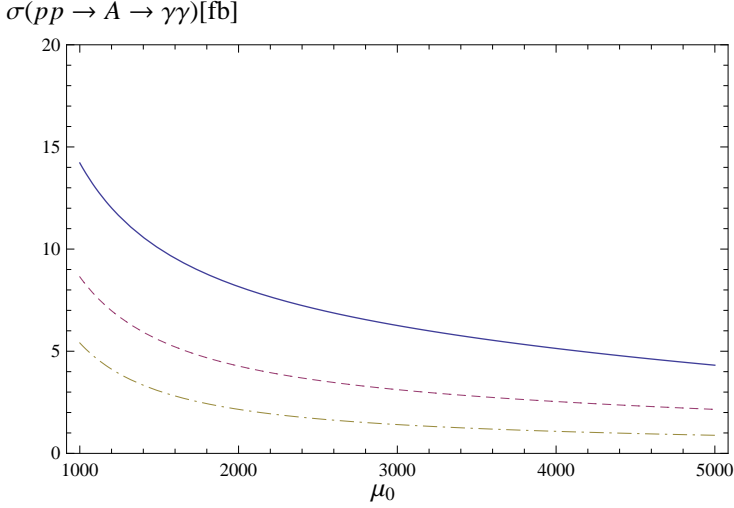


Figure 1. The cross section of $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ is shown as a function of $\mu_Q = \mu_D = \mu_L = \mu_0$ for $\mu_E = 400$ GeV (solid lines), $\mu_E = 500$ GeV (dashed lines) and $\mu_E = 800$ GeV (dashed-dotted lines), for the case $\Lambda_t = 80$ TeV and $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma = 1.5$.

First of all it is worthwhile to identify the scenario that leads to the suppression of the decay rates $A \rightarrow t\bar{t}, WW, ZZ, \gamma Z$ because no indication of the 750 GeV resonance has been observed in the channels associated with these decay modes. The analytical expressions for the corresponding partial decay widths can be presented in the following form

$$\Gamma(A \rightarrow t\bar{t}) = \frac{3m_A m_t^2}{8\pi \Lambda_t^2} \sqrt{1 - \frac{4m_t^2}{m_A^2}}, \quad (15)$$

$$\Gamma(A \rightarrow WW) = \frac{m_A^3}{2\pi} |c_2|^2 \left(1 - \frac{4M_W^2}{m_A^2}\right)^{3/2}, \quad (16)$$

$$\Gamma(A \rightarrow ZZ) = \frac{m_A^3}{4\pi} \left|c_1 \sin^2 \theta_W + c_2 \cos^2 \theta_W\right|^2 \left(1 - \frac{4M_Z^2}{m_A^2}\right)^{3/2}, \quad (17)$$

$$\Gamma(A \rightarrow \gamma Z) = \frac{m_A^3}{8\pi} \sin^2 2\theta_W |c_1 - c_2|^2 \left(1 - \frac{M_Z^2}{m_A^2}\right)^3. \quad (18)$$

In the model under consideration $\Lambda_t \simeq \sqrt{15}f$ if t^c is mainly a component of $\mathbf{20}'$ of SU(6). Since $f \gtrsim 10$ TeV the partial decay width (15) tends to be rather small, i.e. $\Gamma(A \rightarrow t\bar{t}) \lesssim \Gamma(A \rightarrow \gamma\gamma)$.

Here we set $\Lambda_t \simeq 80$ TeV. The partial decay widths (16)–(18) become substantially smaller than $\Gamma(A \rightarrow \gamma\gamma)$ if $|c_2| \ll |c_1|$. The appropriate suppression of $|c_2|$ can be achieved when the exotic fermions that form $SU(2)_W$ doublets are considerably heavier than the $SU(2)_W$ singlet exotic states. On the other hand the non-observation of any new coloured particles with masses below 1 TeV at the LHC implies that the exotic coloured fermions in the E_6 CHM should be rather heavy. Thus to simplify our numerical analysis we assume that $\mu_D = \mu_Q = \mu_L = \mu_0 \gtrsim \mu_E$ and $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma$. For $\mu_0 \gg \mu_E$ the decay rates for $A \rightarrow t\bar{t}, WW, ZZ$ and $Z\gamma$ are very suppressed. However μ_0 cannot be too large, otherwise the LHC production cross section of the pseudoscalar A becomes smaller than 5 – 10 fb. In our numerical analysis we vary μ_0 from 1 TeV to 5 TeV. In this case A mainly decays into a pair of gluons. As a consequence $\Gamma_A \approx \Gamma(A \rightarrow gg)$ and the cross section (13) is determined by $\Gamma(A \rightarrow \gamma\gamma)$. Then the value of μ_E can be adjusted so that $(\Gamma(A \rightarrow \gamma\gamma)/m_A) \sim 10^{-6}$ leading to $\sigma(pp \rightarrow A \rightarrow \gamma\gamma) \simeq 5 - 10$ fb.

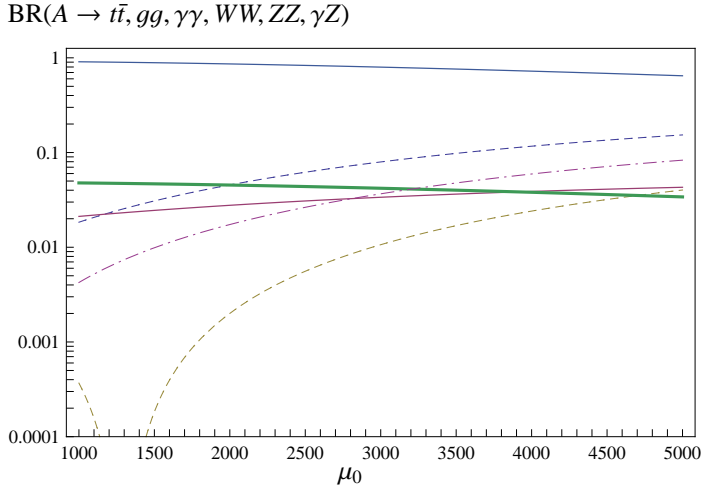


Figure 2. The branching ratios of the decays of the pseudoscalar A into $t\bar{t}$ (dashed–dotted lines), gg (highest solid lines), $\gamma\gamma$ (highest dashed lines), WW (thick solid lines), ZZ (lowest solid lines) and γZ (lowest dashed lines) are presented as a function of $\mu_Q = \mu_D = \mu_L = \mu_0$ for $\mu_E = 400$ GeV, $\Lambda_t = 80$ TeV and $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma = 1.5$.

Fig. 1 demonstrates that $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ decreases very substantially when μ_E increases. For $\mu_E = 400$ GeV and $\sigma \gtrsim 1.5$ the cross section (16) can be about of 5 fb, even when all other exotic fermions are rather heavy $\mu_0 \simeq 5$ TeV. At the same time for $\mu_E \simeq 700 - 800$ GeV the corresponding diphoton production cross section becomes sufficiently large only if $\mu_0 \simeq 1$ TeV. When μ_0 changes from 5 TeV to 1 TeV the ratio Γ_A/m_A increases from 10^{-5} to 10^{-4} that corresponds to the variation of the total LHC production cross section of the pseudoscalar A from 100 fb to 1 pb.

In Fig. 2 the dependence of the branching ratios of the pseudoscalar A on μ_0 is explored for $\mu_E \simeq 400$ GeV and $\sigma = 1.5$. From this figure it follows that $A \rightarrow gg$ is the dominant decay channel. Its branching fraction is always close to 100%. When $\mu_0 \simeq 5$ TeV the branching ratio $BR(A \rightarrow \gamma\gamma)$ is the second largest one. $BR(A \rightarrow WW)$, $BR(A \rightarrow ZZ)$, $BR(A \rightarrow Z\gamma)$ and $BR(A \rightarrow t\bar{t})$ are substantially smaller than $BR(A \rightarrow \gamma\gamma)$. This might be a reason why the decays $A \rightarrow WW, ZZ, Z\gamma, t\bar{t}$ have not been detected yet. The decays $A \rightarrow gg$ can be rather problematic to observe because the total LHC production cross section of the pseudoscalar A is quite small. The branching fractions $BR(A \rightarrow ZZ)$ and $BR(A \rightarrow WW)$ increase while $BR(A \rightarrow \gamma\gamma)$ and $BR(A \rightarrow t\bar{t})$ decrease with decreasing μ_0 .

When $\mu_0 \simeq 1$ TeV the branching ratios $\text{BR}(A \rightarrow ZZ)$ and $\text{BR}(A \rightarrow WW)$ are somewhat bigger than $\text{BR}(A \rightarrow \gamma\gamma)$. Nevertheless the experimental detection of $A \rightarrow ZZ$ and $A \rightarrow WW$ can be rather difficult since the W and Z bosons decay predominantly into quarks. In this scenario $\text{BR}(A \rightarrow \gamma Z)$ remains the lowest branching fraction and vanishes at some value of μ_0 .

This work was supported by the University of Adelaide and the Australian Research Council through the ARC Center of Excellence in Particle Physics at the Terascale (CE 110001004) and through grant LF099 2247 (AWT).

References

- [1] K. Agashe, R. Contino, A. Pomarol, Nucl. Phys. B **719**, 165 (2005).
- [2] B. Bellazzini, C. Csáki, J. Serra, Eur. Phys. J. C **74**, 2766 (2014).
- [3] R. Nevzorov, A. W. Thomas, Phys. Rev. D **92**, 075007 (2015).
- [4] S. F. King, S. Moretti, R. Nevzorov, Phys. Rev. D **73**, 035009 (2006).
- [5] S. F. King, S. Moretti, R. Nevzorov, Phys. Lett. B **634**, 278 (2006).
- [6] S. F. King, S. Moretti, R. Nevzorov, AIP Conf. Proc. **881**, 138 (2007).
- [7] S. F. King, S. Moretti, R. Nevzorov, Phys. Lett. B **650**, 57 (2007).
- [8] S. F. King, R. Luo, D. J. Miller, R. Nevzorov, JHEP **0812**, 042 (2008).
- [9] P. Athron, S. F. King, D. J. Miller, S. Moretti, R. Nevzorov, Phys. Lett. B **681**, 448 (2009).
- [10] P. Athron, S. F. King, D. J. Miller, S. Moretti, R. Nevzorov, Phys. Rev. D **80**, 035009 (2009).
- [11] J. P. Hall, S. F. King, R. Nevzorov, S. Pakvasa, M. Sher, Phys. Rev. D **83**, 075013 (2011).
- [12] P. Athron, S. F. King, D. J. Miller, S. Moretti, R. Nevzorov, Phys. Rev. D **84**, 055006 (2011).
- [13] P. Athron, S. F. King, D. J. Miller, S. Moretti, R. Nevzorov, Phys. Rev. D **86**, 095003 (2012).
- [14] R. Nevzorov, Phys. Rev. D **87**, 015029 (2013).
- [15] R. Nevzorov, S. Pakvasa, Phys. Lett. B **728**, 210 (2014).
- [16] R. Nevzorov, Phys. Rev. D **89**, 055010 (2014).
- [17] P. Athron, M. Mühlleitner, R. Nevzorov, A. G. Williams, JHEP **1501**, 153 (2015).
- [18] P. Athron, D. Harries, R. Nevzorov, A. G. Williams, Phys. Lett. B **760**, 19 (2016).
- [19] S. F. King, R. Nevzorov, JHEP **1603**, 139 (2016).
- [20] M. Frigerio, J. Serra, A. Varagnolo, JHEP **1106**, 029 (2011).
- [21] J. Barnard, T. Gherghetta, T. S. Ray, A. Spray, JHEP **1501**, 067 (2015).
- [22] K. Agashe, R. Contino, R. Sundrum, Phys. Rev. Lett. **95**, 171804 (2005).
- [23] J. Rich, M. Spiro, J. Lloyd–Owen, Phys. Rept. **151**, 239 (1987).
- [24] T. K. Hemmick et al., Phys. Rev. D **41**, 2074 (1990).
- [25] R. Nevzorov, A. W. Thomas, arXiv:1605.07313 [hep-ph].
- [26] R. Franceschini et al., JHEP **1603**, 144 (2016).